Braaten-Pisarski Method at Finite Chemical Potential [*]

Hans Vija[†] and Markus H. Thoma

Institut für Theoretische Physik, Universität Giessen,

35392 Giessen, Germany[‡]

Abstract

The effective perturbation theory developed by Braaten and Pisarski for gauge theories at finite temperature is extended to finite chemical potential. As a first application the collisional energy loss of a heavy quark propagating through a quark-gluon plasma with non-vanishing quark chemical potential is considered. Assuming $\mu/T \simeq 1$, motivated by numerical simulations of heavy ion collisions at RHIC energies, we find that the effect of the quark chemical potential is rather small, unless the energy density instead of the temperature is fixed.

I. INTRODUCTION

QCD lattice calculations predict a phase transition from hadronic matter to a phase of deconfined quarks and gluons (QGP) above a critical energy density $\epsilon_c \approx 1$ GeV/fm^3 corresponding to a critical temperature $T_c \approx 150$ MeV (see, e.g. [1]). The creation of a QGP is expected to happen in ultrarelativistic heavy ion collision and in the early universe shortly after the big bang. An overview of the present status of the experiments is given in [2].

Depending on the available center of mass energy of the colliding nuclei, one can think of two scenarios [3]: Below an energy per nucleon of approximately 5 GeV/A complete stopping of the participating nuclei takes place. It is believed that in this region of high compression and high temperature a QGP containing more quarks than anti-quarks might be produced. Above an energy of about a few 100 GeV/A, on the other hand, full transparency is expected. The QGP is believed to be created in the region between the two expelling and highly excited nuclei by vacuum polarization leading to a vanishing baryon number in the fireball.

We consider the quark chemical potential μ, which is related to the baryo-chemical potential μ_b through $\mu = \mu_b/3$ and measures the deviation from the balance of quarks and anti-quarks. An equal amount of quarks and anti-quarks corresponds to $\mu = 0$, whereas $\mu > 0$ describes an excess of quarks over anti-quarks. It is generally agreed that $\mu/T \simeq 0$ at LHC energies. At RHIC energies, however, there might still exist a considerable amount of stopping leading to $\mu/T \simeq 1$ - 2 [4] as predicted by numerical simulations (RQMD) [5].

Properties of a QGP at finite chemical potential cannot be described by lattice calculations so far. Perturbative QCD at finite temperature and chemical potential suffers from infrared singularities and gauge dependent results. These problems can be avoided for the most physical quantities of interest by adopting an effective perturbation theory developed by Braaten and Pisarski, which is based on the use of resumed Green functions instead of bare ones [6]. In this way consistent results, i.e. complete to leading order in the coupling constant and gauge independent, can be derived in the weak coupling limit, $g \ll 1$.

Here we will consider the Braaten-Pisarski method at finite temperature as well as chemical potential, which is relevant for ultrarelativistic heavy ion collisions if no complete transparency occurs. Up to now only the cases $T \neq 0$, $\mu = 0$ [6] and $T = 0$, $\mu \neq 0$ in an astrophysical context [7] have been discussed.

The next section treats the technical aspects of introducing the chemical potential into thermal field theory. We extend the Braaten-Pisarski method to finite chemical potential and apply the results to the calculation of the energy loss of a heavy quark in a QGP as a typical example in the third section. Throughout the paper we use the metric tensor $g = (+ - --)$, natural units ($\hbar = c = k = 1$) and the notation $K = (k_0, \mathbf{k})$ and $k = |\mathbf{k}|$ for the momenta.

II. PERTURBATIVE TECHNIQUES AT FINITE CHEMICAL POTENTIAL

A. Saclay method

The Braaten-Pisarski method starts from the imaginary time formalism where one has to sum over discrete energies (Matsubara frequencies). For this purpose the Saclay method was developed in the sixties by Balian and De Dominics [8] and used by Pisarski in [9] allowing an elegant way to evaluate finite temperature loops. The idea is based upon the 'mixed' representation of the propagators, i.e. a momentum representation of space components and a space representation of the time component. In the case of the bare fermion propagator including the chemical potential and neglecting the bare mass it reads:

$$S(K) = \frac{1}{\rlap{/}{K}} = \rlap{/}{K} \, \widetilde{\Delta}(K),$$

$$\widetilde{\Delta}(K) \equiv \frac{1}{K^2} = - \int_0^\beta d\tau \, e^{k_0 \tau} \, \widetilde{\Delta}(\tau, k),$$

$$\widetilde{\Delta}(\tau, k) = -T \sum_{n=-\infty}^{\infty} e^{-k_0 \tau} \, \frac{1}{k_0^2 - k^2},$$

$$k_0 = (2n + 1) \, i \, \pi \, T + \mu. \tag{1}$$

The sum is evaluated best by a complex integration [10] yielding the Saclay representation of the fermion propagator:

$$\widetilde{\Delta}(\tau, k) = \frac{1}{2k} \left\{ e^{-k\tau} \left[1 - n_F(k - \mu)\right] - e^{k\tau} n_F(k + \mu) \right\}, \tag{2}$$

where $n_F(k \pm \mu) = 1/[\exp((k \pm \mu)/T) + 1]$ denotes the Fermi distribution function. The advantage of this representation is the easy summation over the Matsubara frequencies since the discrete energy k_0 appears only in $\exp(k_0 \tau)$ in the propagator $\widetilde{\Delta}(K)$. The derivation of the boson propagator containing no chemical potential proceeds analogously [9].

B. Hard Thermal Loops

Following the program given by the Braaten-Pisarski resummation technique in ref. [6] we calculate the Hard Thermal Loops (HTL) at finite chemical potential, first. HTL are one-loop diagrams for which only the contribution from hard loop momenta of the order T or larger is considered. This approximation is equivalent to the leading term of the high temperature expansion of the diagram under consideration. Here, we attach importance mainly to changes introduced by the chemical potential and refer to the literature otherwise.

First we discuss the HTL approximation of the gluon self energy, including the quark chemical potential and using the Saclay method. The only modification to the HTL polarization tensor compared to the zero chemical potential case comes from the diagram of fig.1. Using standard Feynman rules in the imaginary time formalism and the Saclay representation of the propagators, where we consider one fermion line as a quark with chemical potential μ and the other one as an anti-quark with $-\mu$, we obtain analogously to the case of zero chemical potential [6], assuming the HTL limit,

$$\Pi_L(K) \equiv \Pi_{00}(K)$$
$$= -3\, m_g^2 \left(1 - \frac{k_0}{2k} \ln \frac{k_0 + k}{k_0 - k}\right) \tag{3}$$

for the longitudinal part and

$$\Pi_T(K) \equiv \frac{1}{2} \left(\delta_{ij} - \frac{k_i k_j}{k^2} \right) \Pi_{ij}$$

$$= \frac{3}{2} m_g^2 \frac{k_0^2}{k^2} \left[1 - \left(1 - \frac{k^2}{k_0^2} \right) \frac{k_0}{2k} \ln \frac{k_0 + k}{k_0 - k} \right] \qquad (4)$$

for the transverse part of the gluon polarization tensor. Here

$$m_g^2 = \frac{g^2 T^2}{3} \left(1 + \frac{\tilde{N}_f}{6} \right),$$

$$\tilde{N}_f = N_f + \frac{3}{\pi^2} \sum_f \frac{\mu_f^2}{T^2}, \qquad (5)$$

is called the effective gluon mass, where N_f denotes the number of thermal flavors in the QGP and μ_f the corresponding chemical potentials. In the following a QGP consisting of up and down quarks with $\mu_u = \mu_d = \mu$ and all other $\mu_f = 0$ is assumed. The only change in the gluon self energy due to the introduction of a chemical potential is given by the effective gluon mass which depends now on the chemical potential in accordance with the result of the high temperature limit [11]. In the limit $T \to 0$ the result (3) to (5) also reduces to the one found in ref. [7].

The same behavior is observed for the HTL quark self energy. One simply has to replace the effective quark mass by the chemical potential dependent one

$$m_q^2 = \frac{g^2 T^2}{6} \left(1 + \frac{1}{\pi^2} \frac{\mu^2}{T^2} \right) \qquad (6)$$

already quoted in ref. [12]. (In the course of the derivation of the HTL quark self energy we had to assume that the chemical potential is not much larger than the temperature, $\mu \ll T/g$, whereas this restriction was not necessary for the gluon self energy.)

Owing to Ward identities the HTL self energies also imply the existence of HTL vertex corrections [6]. They also will be modified in the presence of a finite chemical potential in a similar way as the self energies. However, for the most applications considered so far they can be neglected.

C. Effective propagators

The effective gluon and quark propagators follow by resuming the HTL self energies by means of a Dyson-Schwinger equation [6]. Hence they differ from the ones at zero chemical potential (see, e.g. ref. [13] for the gluon propagator and ref. [14] for the quark propagator) only by the effective masses (5) and (6). According to the Braaten-Pisarski method effective Green functions (propagators or vertices) have to be used only if all legs of the Green function have soft momenta of the order gT. Otherwise the use of bare Green functions is sufficient. In this way gauge independent results, which are complete to leading order in the coupling constant, are obtained for physical quantities. Moreover, due to the self energies in the denominator of the effective propagators important medium effects such as Debye screening, which improves the infrared behavior, and Landau damping are included. Thus the Braaten-Pisarski method means a crucial improvement compared to the naive perturbation theory based on bare Green functions. However, the application of the Braaten-Pisarski method is limited by two problems [15]. First, the Braaten-Pisarski method relies on the distinction between hard and soft momenta, i.e. on the weak coupling limit assumption, $g \ll 1$. The validity of the extrapolation of the results to realistic values of the coupling constant, $g = 1.5$ - 2.5, in ultrarelativistic heavy ion collisions, has to be checked for each quantity individually. In general it appears to work for particles with large momenta, $p \gg T$. Secondly, the effective gluon propagator suffers from the absence of a static magnetic screening mass in the HTL approximation of the gluon polarization tensor, thus still leading to infrared singularities in certain quantities, as e.g. the damping rate of a hard parton [16]. However, so far only a few observables, e.g. the color relaxation [17], are known to be affected by this problem. The most physical quantities of interest, such as the collisional energy loss of energetic partons [18], transport rates of quarks and gluons [19], the viscosity of the QGP [20], and the production rate of hard photons in the QGP [21], are to leading order infrared finite by using effective propagators due to dynamical magnetic screening [22].

Summarizing this section, the generalization of the Braaten-Pisarski method to finite

chemical potential is straightforward by modifying the effective parton masses, containing now also the chemical potential, that enter into the effective propagators.

III. ENERGY LOSS OF A HEAVY QUARK

The energy loss of energetic partons in a QGP is related to the phenomenon of jet quenching in ultrarelativistic heavy ion collisions [23], which might serve as a signature for the QGP formation [24]. The energy loss of a parton in a QGP is caused by elastic scattering off the thermal partons (collisional energy loss) and gluon bremsstrahlung (radiative energy loss). The collisional contribution follows from the damping rate by introducing the energy transfer per parton collision as a weight under the integral defining the damping rate [13].

The damping rate of a hard quark follows from the imaginary part of the quark self energy diagram shown in fig.2, which comes from the imaginary part of the effective gluon propagator corresponding to Landau damping of a virtual collective gluon mode. This damping mechanism is equivalent to elastic scattering off the thermal partons via the exchange of a collective gluon (plasma wave), as can be seen from cutting the diagram in fig.2. As mentioned above, the damping rate, which is quadratic infrared divergent using bare propagators, turns out to be logarithmically infrared divergent applying the Braaten-Pisarski method reflecting the absence of static magnetic screening in the effective gluon propagator. Integrating over the loop momentum a cancellation of the effective gluon mass takes place [19]. Consequently the damping rate of a hard quark, given to logarithmic accuracy by $\gamma_q = g^2 T/(3\pi) \ln(1/g)$ [25], does not depend on the chemical potential. (The Fermi-Dirac distributions from the internal quark line is exponentially suppressed for hard quarks [13].)

The additional energy transfer factor in the loop integral reduces the infrared divergence and leads to a finite result using the effective gluon propagator at soft momentum transfer, $q \lesssim gT$. For large momentum transfer, $q \gtrsim T$, the exchange of a bare gluon in the elastic scattering diagram is sufficient. According to the prescription by Braaten and Yuan [26] the both contributions can be matched onto each other by assuming a separation scale,

$gT \ll q^\star \ll T$, in the weak coupling limit. Adding up the soft and hard contributions, the separation scale drops out of the final result. In this way a infrared finite and gauge independent result for the collisional energy loss of a heavy quark is found, which is complete to leading order of the coupling constant and independent of arbitrary parameters such as infrared and ultraviolet momentum cut-offs or separation scales.

Taking a finite quark chemical potential into account, the derivation of the soft part proceeds analogously to the case of vanishing chemical potential [13,18]. The only change comes from the effective gluon mass that does not cancel as opposed to the case of the damping rate. The hard part requires integrations over Fermi-Dirac distributions, which can be achieved only numerically at finite chemical potential. It involves integrals of the form

$$\int_0^\infty dk \, [n_F(k - \mu) + n_F(k + \mu)] \, k = \frac{\pi^2 T^2}{6} \left(1 + \frac{3\mu^2}{\pi^2 T^2}\right),$$

$$\int_0^\infty dk \, [n_F(k - \mu) + n_F(k + \mu)] \frac{k}{q^\star} = \frac{\pi^2 T^2}{6} \left(1 + \frac{3\mu^2}{\pi^2 T^2}\right)$$

$$\left[\ln \frac{T}{q^\star} + \ln 2 + 1 - \gamma + \frac{\zeta'(2)}{\zeta(2)} + F\left(\frac{\mu}{t}\right)\right], \qquad (7)$$

where $\gamma = 0.57722...$ is Euler's constant, $\zeta(z)$ Riemann's zeta function with $\zeta'(2)/\zeta(2) = -0.56996...$, and the function $F(\mu/T)$ with $F(0) = 0$, shown in fig.3, has to be determined numerically.

The final result for the energy loss of a heavy quark with mass $M \gg T$, energy E, and velocity v in a QGP of N_f thermal flavors after adding up the soft and hard contributions is given by rather simple expressions in the limits $E \ll M^2/T$,

$$-\frac{dE}{dx} = \frac{8\pi \alpha_s^2 T^2}{3} \left(1 + \frac{\tilde{N}_f}{6}\right) \left(\frac{1}{v} - \frac{1 - v^2}{2v^2} \ln \frac{1 + v}{1 - v}\right)$$

$$\left[\ln \left(2^{\frac{\tilde{N}_f}{6 + \tilde{N}_f}} B(v) \frac{ET}{m_g M}\right) + \frac{\tilde{N}_f}{6 + \tilde{N}_f} F\left(\frac{\mu}{T}\right)\right], \qquad (8)$$

and $E \gg M^2/T$,

$$-\frac{dE}{dx} = \frac{8\pi \alpha_s^2 T^2}{3} \left(1 + \frac{\tilde{N}_f}{6}\right) \left[\ln \left(2^{\frac{\tilde{N}_f}{2(6 + \tilde{N}_f)}} 0.920 \frac{\sqrt{ET}}{m_g}\right) + \frac{\tilde{N}_f}{2(6 + \tilde{N}_f)} F\left(\frac{\mu}{T}\right)\right] \qquad (9)$$

8

respectively, where $B(v)$ is a smooth function of the velocity that increases monotonically from $B(0) = 0.604$ to a maximum of $B(0.88) = 0.731$, and then decreases to $B(1) = 0.629$ [18]. The expressions (8) and (9) extend the Bethe-Bloch formula [27] to the case of a relativistic plasma with finite temperature and chemical potential. They differ from the corresponding expressions at zero chemical potential [18] by the effective gluon mass (5), the effective flavor number $\tilde{N}_f$, and the additional function $F(\mu/T)$. The details of the derivation of (8) and (9) are presented in ref. [28].

Figs.4 and 5 show the energy loss of a charm and bottom quark extrapolated to $\alpha_s = 0.2$ at $T = 250$ MeV and $N_f = 2$ for $\mu/T = 0$, 1, and 2 as a function of the momentum p. In the case of a charm quark (8) is matched onto (9) at some intermediate momentum by demanding dE/dx to be continuous, whereas for a bottom quark (8) is sufficient for the momentum range chosen in fig.5. At momenta below a few GeV the energy loss turns out to be negative reflecting the breakdown of the extrapolation to realistic values of the coupling constant for small momenta [15,18]. The enhancement of dE/dx with increasing chemical potential is rather small (below about 25% at $\mu/T = 2$) since there is a partial cancellation between the factor $\tilde{N}_f$ increasing with μ and the function $F(\mu/T)$ which is negative for $\mu > 0$ (fig.3). This demonstrates that the temperature is more important than the chemical potential for the energy loss.

However, if one fixes the energy density, given e.g. by the MIT-bag equation of state [29]

$$\epsilon(\mu, T) = \left(\frac{37\pi^2}{30} - \frac{11\pi\alpha_s}{3}\right) T^4 + 3\left(1 - \frac{2\alpha_s}{\pi}\right) T^2 \mu^2 + \frac{3}{2\pi^2}\left(1 - \frac{2\alpha_s}{\pi}\right) \mu^4 + B \qquad (10)$$

with the bag constant $B = (200 \text{ MeV}^4)$, increasing the chemical potential corresponds to a decrease of the temperature resulting in a significant reduction of the energy loss as shown in fig.6. (At $\mu = 0$ a temperature of $T = 250$ MeV corresponds to $\epsilon = 5.232$ GeV/fm^3.) Considering ultrarelativistic heavy ion collisions the choice of a certain energy density instead of a temperature might be reasonable as the former is closer related to the collision energy [30].

The collisional energy loss of a light quark follows from (9) by replacing the energy E by a

maximum momentum transfer $\tilde{q} \simeq E/2$, the energy loss of a gluon by multiplying the one of a light quark by a color factor of 9/4 [31]. The radiative energy loss by gluon bremsstrahlung is much more complicated and has not been addressed by the Braaten-Pisarski method so far. Estimates based on naive perturbation theory indicate the dominance of the radiative energy loss over the collisional one in a QGP [32].

IV. CONCLUSIONS

In the present work the Braaten-Pisarski method has been extended to finite chemical potential. The effective propagators, based on a resummation of the HTL self energies, differ from the one at zero chemical potential only by the effective masses which now depend on the chemical potential. The Braaten-Pisarski method can be used for calculating consistently quantities of the QGP that are logarithmically infrared divergent if bare Green Functions are utilized. Interesting observables related to the evolution and detection of a QGP in ultrarelativistic heavy ion collisions, belonging to this class of quantities, are the energy loss of energetic partons propagating through the QGP, the transport rates of quarks and gluons, the viscosity of the QGP, and the production rate of hard photons in the QGP.

Here the collisional energy loss of a heavy quark, representing a typical example for the application of the Braaten-Pisarski method, has been computed to leading order at finite temperature as well as finite quark chemical potential. This might be relevant for heavy ion collisions at RHIC energies, for which numerical simulations (RQMD) predict a considerable amount of stopping leading to a quark chemical potential of the order of $\mu \simeq 1$ - $2 \, T$. The energy loss of a heavy quark has been calculated consistently adopting the Braaten-Yuan prescription, which corresponds to a decomposition in a soft contribution, calculable with an effective gluon propagator, and a hard, within naive perturbation theory accessible contribution. Besides the modification caused by the effective gluon mass (5) in the soft part, momentum integrations over the Fermi-Dirac distribution functions in the hard part lead to an additional term (see (7)), which requires a numerical evaluation at finite

10

chemical potential. However, these changes are small, typical below 25%, for $\mu \leq 2T$ as long as the temperature is fixed, showing the dominant role of the temperature over the chemical potential. If, on the other hand, a given energy density $\epsilon(\mu, T)$ is assumed, corresponding to a decrease of the temperature with increasing chemical potential, a significant suppression is observed.

The same conclusions apply to the transport rates of quarks and gluons, which determine thermalization times and the viscosity of the QGP [19], The transport rate follows from the damping rate by including a transport weight describing momentum relaxation. Owing to this transport weight the computation of the transport rate proceeds similarly to the one of the energy loss via the Braaten-Yuan prescription. The details of this calculation can be found in ref. [28].

The production rate of hard photons from a QGP can be derived to leading order again by using the Braaten-Yuan prescription, where the soft part follows from the photon self energy containing one effective quark propagator and the hard part from the tree matrix elements of Compton scattering and pair annihilation involving one gluon and one photon [21]. The calculation of the hard photon production rate at finite quark chemical potential considering the effective quark propagator modified by the effective quark mass (6) will be presented in a in a forthcoming publication [33].

REFERENCES

[1] B. Petersson, Nucl. Phys. A525 (1991) 237c.

[2] H.R. Schmidt and J. Schuhkraft, J. Phys. G 19 (1993) 1705.

[3] B. Müller, The Physics of the Quark-Gluon Plasma (Springer, Berlin, 1985).

[4] A. Dumitru, T. Schönfeld, D.H. Rischke, H. Stöcker, W. Greiner, and L. McLerran, Frankfurt preprint UFTP 319/1992 (1992).

[5] H. Sorge, H. Stöcker, and W. Greiner, Ann. Phys. (N.Y.) 192 (1989) 266.

[6] E. Braaten and R.D. Pisarski, Nucl. Phys. B337 (1990) 569.

[7] T. Altherr and U. Kraemmer, Astropart. Phys. 1 (1992) 133.

[8] R. Balian and C. De Dominics, Nucl. Phys. 16 (1960) 502.

[9] R. D. Pisarski, Nucl. Phys. B309 (1988) 476.

[10] J.I. Kapusta, Finite Temperature Field Theory (Cambridge University Press, Cambridge, 1989).

[11] T. Toimela, Phys. Lett. B 176 (1986) 463.

[12] K. Kajantie and P.V. Ruuskanen, Phys. Lett. B 121 (1983) 352.

[13] E. Braaten and M.H. Thoma, Phys. Rev. D 44 (1991) 1298.

[14] E. Braaten, R.D. Pisarski, and T.C. Yuan, Phys. Rev. Lett. 64 (1990) 2242.

[15] M.H. Thoma, in: Proc. Workshop on Pre-Equilibrium Parton Dynamics (Berkeley, 1993), ed. X.N. Wang (LBL-34831, 1994) p. 171.

[16] R.D. Pisarski, Phys. Rev. Lett. 63 (1989) 1129; M.H. Thoma and M. Gyulassy, Nucl. Phys. B351 (1991) 491.

[17] A.V. Selikhov and M. Gyulassy, Phys. Lett. B 316 (1993) 373.

[18] E. Braaten and M.H. Thoma, Phys. Rev. D 44 (1991) R2625.

[19] M.H. Thoma, Phys. Rev. D 49 (1994) 451.

[20] M.H. Thoma, Phys. Lett. B 269 (1991) 144.

[21] J.I. Kapusta, P. Lichard, and D. Seibert, Phys. Rev. D 44 (1991) 2774; R. Baier, H. Nakkagawa, A. Niégawa, and K. Redlich, Z. Phys. C 53 (1992) 433.

[22] G. Baym, H. Monien, C.J. Pethick, and D.G. Ravenhall, Phys. Rev. Lett. 64 (1990) 1867.

[23] J.D. Bjorken, Fermilab Pub-82/59-THY (1982), unpublished.

[24] M. Gyulassy and M. Plümer, Phys. Lett. B 243 (1990) 432.

[25] C.P. Burgess and A.L. Marini, Phys. Rev. D 45 (1992) 17.

[26] E. Braaten and T.C. Yuan, Phys. Rev. Lett. 66 (1991) 2183.

[27] J.D. Jackson, Classical Electrodynamics (Wiley, New York, 1975).

[28] H. Vija, diploma thesis, Giessen (1994).

[29] A. Dumitru, D.H. Rischke, H. Stöcker, and W. Greiner, Mod. Phys. Lett. A 8 (1993) 1291.

[30] A. Dumitru, D.H. Rischke, T. Schönfeld, L. Winckelmann, H. Stöcker, and W. Greiner, Phys. Rev. Lett. 70 (1993) 2860.

[31] M.H. Thoma, Phys. Lett. B 273 (1991) 128.

[32] M. Gyulassy and X.N. Wang, Nucl. Phys. B420 (1994) 583.

[33] C. Traxler and M.H. Thoma, work in progress.

FIGURES

FIG. 1. Part of the HTL gluon polarization tensor containing the quark chemical potential

FIG. 2. Quark self energy containing the effective gluon propagator

FIG. 3. The function $F(\mu/T)$

FIG. 4. Energy loss of a charm quark for $\mu/T = 0$, 1 and 2 for a given temperature

FIG. 5. Energy loss of a bottom quark for $\mu/T = 0$, 1 and 2 for a given temperature

FIG. 6. Energy loss of a bottom quark for $\mu/T = 2$ for a given energy density compared to a given temperature

This figure "fig1-1.png" is available in "png" format from:

http://arXiv.org/ps/hep-ph/9409246v1

This figure "fig1-2.png" is available in "png" format from:

http://arXiv.org/ps/hep-ph/9409246v1

This figure "fig1-3.png" is available in "png" format from:

http://arXiv.org/ps/hep-ph/9409246v1